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A b s t r a c t

The electron (positron) polarization process in storage rings is studied when radiating the beam by a colliding electromagnetic wave. The effect arises due to dependence of radiation friction in the wave field on the spin and the spin motion coupling with the orbital motion in the field of a storage ring. The degree of equilibrium polarization and the relaxation time are found. A numerical example for the case of laser polarization is presented.

1. INTRODUCTION

The only realized method of producing the polarized electrons and positrons in storage rings is currently based on the use of a polarizing effect of synchrotron radiation. Efficiency of this method is determined by the particle energy and the magnetic field value of a storage ring ($\tau_{S.R.}^{-1} \sim E^2 H^3$).

Unfortunately, the growth of radiative losses with energy causes a decrease in the leading field value of a storage ring. For this reason, in existing and designed storage rings, despite their high energy, the polarization time remains practically unchangeable over an order of magnitude and equals several hours even at a maximum energy of the storage ring.

Therefore, the search for methods for quickly producing the polarized beams seems to be urgent. The introduction of large magnetic fields (magnetic "snakes") /1-3/ into a relatively small part of the orbit is a known but unrealized method.

In this paper the method of electron polarization is studied which uses a spin dependence of radiation in the field of a colliding circularly polarized electromagnetic wave. The effect may be essential due to high wave frequency.

A polarizing mechanism resulted from the spin dependence of radiation friction force (processes with no spin-flip during radiation) and the orbit-spin coupling in the field of a storage ring was previously found in /4/. It turns out that the polarization by a colliding circularly-polarized wave is possible only in the case of utilizing this mechanism.

The problem of polarizing the electrons (positrons) in a stationary helical magnetic field (in an ondulator) was considered in /5, 6/. The action of the ondulator on the ultrarelativistic electrons is theoretically equivalent to that of a colliding electromagnetic wave. However, the question of a possibility to polarize the particle in a storage ring by such a method was not answered in these papers, since the orbit-spin coupling effects were not borne in mind in them.

2. POLARIZATION KINETICS

Let the electrons move in the straight section of a storage ring on the length ℓ in the field of colliding circularly-polarized wave. We shall consider the case of soft radiation when the particles do not leave the beam because of quantum recoil: $\hbar\omega_{ph} < \Delta E_p$, where ω_{ph} is the characteristic radiation frequency, ΔE_p is the maximally permissible energy jump determined by the aperture of a storage ring. The condition of a quasiclassical character of radiation ($\hbar\omega_{ph} \ll E$) is here fulfilled as well, and it is possible to use the method and the general formulae of /4/. The polarization variation will occur due to multiple passings of the particles through the section with a wave field. Here it should take into account the perturbation of a particle trajectory by radiation. It is convenient to apply a general method commonly used in the theory of orbital-motion rac ative effects. The integrals of motion in the storage ring field, i.e. the integrals of action (or of amplitude) and phase, are firstly determined. Then, the phase-average rates of changes in the action variables under the influence of ra-

diation are found, according to perturbation theory. Because of phase mixing, this suffices to find the relaxation times and the equilibrium state.

In a homogeneous magnetic field the action-phase spin variables are, apparently, the projections of a spin on the field direction and the precession phases around the field. Since in the general case spin motion is a rotation, the action variable is of the form $S_{\vec{n}} = \vec{S}\vec{n}(\vec{P},\vec{z})$, where $\vec{n}$ is a single vector giving the precession axis direction and being a function of the momentum $\vec{P}$ and coordinates $\vec{z}$ of a particle. The precession phase ψ around $\vec{n}$ is canonically conjugated to $S_{\vec{n}}$.

The vector $\vec{n}$ is obviously the time-independent solution of the equation of spin motion in a field of the storage ring and completely determined by a particle trajectory. Correspondingly, the spectrum $\vec{n}$ contains only the frequencies of orbital motion. With respect to spin motion the vector $\vec{n}$ plays the same role as the energy-dependent closed orbits of particles relative to betatron oscillations.

After the mixing over the phases ψ , due to spread in the precession frequencies $\langle\dot{\psi}\rangle$, the mean spin for a group of particles traveling near the equilibrium orbit will be directed along the precession axis on the equilibrium trajectory $\vec{n}_s(\theta) = \vec{n}(\vec{P}_s, \vec{z}_s)$:

$$\langle\vec{S}\rangle = \langle S_{\vec{n}}\vec{n}\rangle + \langle\vec{S}_\perp\rangle = \langle S_{\vec{n}}\rangle\vec{n}_s .$$

The vector $\vec{n}_s$ is, apparently, periodical over the generalized azimuth of the particle θ : $\vec{n}_s(\theta) = \vec{n}_s(\theta+2\pi)$.

The equilibrium polarization is thus directed along $\vec{n}_s$, independent of the polarizing mechanism. Practically, any required direction $\vec{n}_s$ at a given place of the orbit, in particular, on the section with a wave field /2, 7/ may be given in a storage ring.

The polarization degree $\zeta = \langle S_{\vec{n}} \rangle / S$ will be slowly vary under the radiation influence, approaching to some value determined by the equilibrium of polarizing and depolarizing processes.

The mean variation rate $S_{\vec{n}}$ is written as follows:

$$\langle \dot{S}_n \rangle = \left\langle \frac{\delta S_{\vec{n}}}{\delta \tau} \right\rangle = \left\langle \vec{n} \frac{\delta \vec{S}}{\delta \tau} \right\rangle + \left\langle \vec{S} \frac{\delta \vec{n}}{\delta \tau} \right\rangle + 2 \left\langle \frac{\delta \vec{S} \delta \vec{n}}{\delta \tau} \right\rangle , \qquad (1)$$

where $\delta \vec{S}$ and $\delta \vec{n}$ are the increments $\vec{S}$ and $\vec{n}$ under the radiation influence at time $\delta \tau$. The first term $\langle \vec{n} \delta \vec{S} / \delta \tau \rangle$ describes the direct radiation effect on the spin, the second - the radiation effect on the precession axis by the orbital motion perturbation, the last takes into account the correlation between the jumps of the spin and the particle momentum in quantum radiation.

The change in the precession axis $\vec{n}$ is expressed through the change in a particle momentum $\vec{P}$ in radiation:

$$\delta \vec{n} \approx \sum_{\alpha=1}^{3} \frac{\partial \vec{n}}{\partial P_\alpha} \delta P_\alpha + \frac{1}{2} \sum_{\alpha,\beta=1}^{3} \frac{\partial^2 \vec{n}}{\partial P_\alpha \partial P_\beta} \delta P_\alpha \delta P_\beta . \qquad (2)$$

In the general case, Eq.(1) is reduced to the form:

$$\dot{\zeta} = -\alpha_+ \zeta + \alpha_- , \qquad (3)$$

wherein the coefficients $\alpha_\pm$ denote, correspondingly, the sum and differencie of the spin-flip probabilities per unit time relative to the vector direction $\vec{n}$. The degree of polarization exponentially tends to the value $\zeta_\infty = \alpha_- / \alpha_+$, and the relaxation time equals $\tau_p = \alpha_+^{-1}$. The general expression of the coefficients $\alpha_\pm$ in the integral form over classical trajec-

tories is presented in /4/ (formulae (4.7)). In the case of high wave frequency,[1] $\omega_w \gg \frac{e}{m} H_w$, which is of interest to us, the calculations lead to the following results (the influence of synchrotron radiation on the sections with a leading field is ignored):

$$\alpha_- = -8\hbar q_0^4 \zeta_2 \omega_w \gamma^2 H_w^2 \left\{ \frac{2}{3} \gamma \vec{v} \frac{\partial \vec{n}}{\partial \gamma} + \frac{m}{6} \sum_{\alpha=1}^{2} \frac{\partial n_\alpha}{\partial P_\alpha} \right\} \frac{\ell}{L} ,$$

$$\alpha_+ = 8\hbar q_0^4 \omega_w \gamma^2 H_w^2 \left\{ \frac{2}{3} - \frac{4}{15} (\vec{n}\vec{v})^2 + \frac{7}{15} \left(\gamma \frac{\partial \vec{n}}{\partial \gamma} \right)^2 - \frac{m}{30} \vec{n} \left[\left[\vec{v} \frac{\partial}{\partial \vec{p}} \right] \vec{n} \right] + \frac{m^2}{10} \sum_{\alpha=1}^{2} \left(\frac{\partial \vec{n}}{\partial P_\alpha} \right)^2 \right\} \frac{\ell}{L} , \qquad (4)$$

where $q_0 = \frac{e}{m}$, $(c=1)$, H_w is the magnetic wave field in the laboratory system, $\gamma = (1-v^2)^{-1/2}$ is a relativistic factor $(= P_3/m)$, $P_{1,2}$ are the transverse components of momentum, L is the perimeter of a storage ring, ζ_2 is the circular polarization degree of the wave which, particularly, equals ± 1 for the right- and left-polarized wave, respectively.

In similar expressions for α_+ and α_- the factor 8 is absent in the ondulator case.

The coefficients $\alpha_\pm$ may be also calculated by means of the known cross-section of Compton scattering /8/. With this method of calculation the variation in a direction $\vec{n}$ after scattering is also taken into account and is of essential importance.

The Compton scattering cross-section of the electron on a colliding wave which are summed over polarizations of a finite

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1) For the opposite case, when $\omega_w \ll \frac{e}{m} H_w$ (the case of synchrotron radiation) the conclusive answer is in /4/.

state of the photon may be written as a function of electron-spin orientations before and after scattering $\vec{\zeta}$ and $\vec{\zeta}'$, and of the variation in the momentum $\delta\vec{P}$:

$$d\sigma = d\sigma(\vec{\zeta}, \vec{\zeta}'; \delta\vec{P}).$$

Hence, flip probabilities $P_{\uparrow\downarrow}$ and $P_{\downarrow\uparrow}$ with respect to the variable direction $\vec{n}$ which determine the coefficients $\alpha_\pm$ $(\alpha_\pm = P_{\uparrow\downarrow} \pm P_{\downarrow\uparrow})$ may be found from the formulae

$$P_{\uparrow\downarrow} = \dot{N}_W \int d\sigma(\vec{n}, -\vec{n}'; \delta\vec{P}),$$
$$P_{\downarrow\uparrow} = \dot{N}_W \int d\sigma(-\vec{n}, \vec{n}'; \delta\vec{P}), \qquad (5)$$

where the photon flux density is denoted as $\dot{N}_W$ $(\dot{N}_W = H_W^2/(4\pi\hbar\omega_W))$, and $\vec{n}' = \vec{n} + \delta\vec{n}$ is defined from formula (2). In calculation of the probabilities in the expression for $d\sigma$ one needs to remain the terms up to the order $(\hbar\omega_{ph}/E)^2$.

In the expression for the scattering cross-section the value $\partial^2\vec{n}/\partial P_\alpha \partial P_\beta$ is included in the form of projection on the direction $\vec{n}$. This enables one to express the result through the values $\partial\vec{n}/\partial P_\alpha$ $(\vec{n}^2 = 1)$:

$$\vec{n}\frac{\partial^2\vec{n}}{\partial P_\alpha \partial P_\beta} = -\frac{\partial\vec{n}}{\partial P_\alpha}\frac{\partial\vec{n}}{\partial P_\beta}.$$

In order to use the expression for the Compton scattering cross-section in an accompanying system, it is necessary to express the spin direction before and after scattering in the accompanying system through the corresponding directions in the laboratory system. Since the electron in the accompanying system before scattering rests, the directions $\vec{n}$ and $\vec{n}_a$ coincide. After scattering the polarization directions in these two systems will differ by Thomas revolution because of the variation in momentum. In the linear approximation over δP we have:

$$\vec{n}_a = \vec{n}' + \frac{\gamma - 1}{\gamma m v^2}[[\vec{v}\delta\vec{P}]\vec{n}].$$

It is easy to find from this a connection between the quantity $(\partial\vec{n}_a/\partial P_\alpha)_a$ in the accompanying system and its value in the laboratory system $(\gamma \gg 1)$:

$$\left(\frac{\partial\vec{n}_a}{\partial P_\alpha}\right)_a = \frac{\partial\vec{n}}{\partial P_\alpha} + \frac{1}{m}[[\vec{v}\vec{e}_\alpha]\vec{n}],$$
$$m\left(\frac{\partial\vec{n}_a}{\partial P_3}\right)_a = \gamma\frac{\partial\vec{n}}{\partial\gamma}. \qquad (6)$$

3. DISCUSSION OF RESULTS

Explain now a physical sense of each term in the expressions for α_+ and α_-.

Let us consider the first term $\langle\vec{n}\frac{\delta\vec{S}}{\delta\tau}\rangle$ in the kinetic equation (1). As has been already mentioned above, it describes the direct radiation effect on a spin, i.e. the scattering of photons with electron spin flip. The appearance of the terms independent of the gradients $\partial\vec{n}$ in expressions (4) for $\alpha_\pm$ is due to just these scattering processes.

Supposing that $\vec{n} = \vec{n}'$ in formulae (5), it is possible to find the flip probabilities $P^s_{\uparrow\downarrow}$, $P^s_{\downarrow\uparrow}$ describing the direct radiation effect on a spin via the Compton scattering cross-section:

$$P^s_{\uparrow\downarrow} = \dot{N}_W \int d\sigma(\vec{n}, -\vec{n}; \delta\vec{P}),$$
$$P^s_{\downarrow\uparrow} = \dot{N}_W \int d\sigma(-\vec{n}, \vec{n}; \delta\vec{P}).$$

In the ultrarelativistic limit these probabilities seems to be equal and, hence, the direct radiation effect is depolarizing.

The second term in (1) takes into account the effect of the variation in a particle momentum in scattering on the precession axis. The polarization effect here arises due to correlation of the direction of spin $\vec{S}$ and that of the particle momentum during scattering without a spin flip. This polarizing effect is described by the expression of the form:

$$\langle \dot{S}_{\vec{n}} \rangle_{pol} = \langle \vec{S} \frac{\delta \vec{n}}{\delta \tau} \rangle_{pol} = \langle \vec{S} (\vec{f}_s \frac{\partial}{\partial \vec{P}}) \vec{n} \rangle , \qquad (7)$$

where $\vec{f}_s$ is the part of the radiative friction force which is linear over the spin $\vec{S}$. The latter is of the form:

$$\vec{f} = -\frac{8}{3} q_0^2 \gamma^2 H_w^2 \vec{v} + \frac{8}{3} m q_0^4 \gamma^3 \zeta_2 \omega_w H_w^2 \left[4(\vec{S}\vec{v})\vec{v} + \frac{1}{\gamma} \vec{S}_\perp \right]$$

$$(\vec{S}_\perp = \vec{S} - (\vec{S}\vec{v})\vec{v}/v^2)$$

and may be obtained by averaging the change in a particle momentum over the scattering probability $\dot{N}_w d\sigma$:

$$\vec{f} = \dot{N}_W \int \delta \vec{P} \, d\sigma (\vec{\zeta}, \vec{\zeta}'; \delta \vec{P}) .$$

The force $\vec{f}$ as a function of the spin $\vec{S} = \frac{\hbar}{2} \vec{\zeta}$ does not contain the Plank constant and may be obtained in the classical radiation theory for a particle with magnetic moment and spin. Substitution[2)] of the force $\vec{f}$ into (7) leads to the result consistent with the expression for α_-.

Observin_ the dynamics of spin and orbital motion, one can explain directly the polarizing effect described by a parameter $\partial \vec{n} / \partial P_\alpha$. After a single passing of the interaction region

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2) In quantum consideration of some terms in (1) they should be replaced by the spin-symmetrized expressions, for example: $\hat{S}_\alpha f_\beta \to \frac{1}{2} (\hat{S}_\alpha \hat{f}_\beta + \hat{f}_\beta \hat{S}_\alpha)$.

the spin orientation does not vary, but under the radiation influence the electron receives the momentum whose average value depends on a spin direction. The orbital perturbation in an inhomogeneous field of the storage ring results in varying the field affecting the spin and correlated with it. At synchrotron radiation this perturbation relaxes at a larger number of turns in the storage ring, thereby leading to the spin-quadratized integral effect depending upon a relative position of the spin and orbital motion frequencies.

The quantum radiation fluctuations in a wave field leads to the diffusion change in a particle momentum which in turn results in the diffusion change of the projection $S_{\vec{n}}$ in the case when the precession axis $\vec{n}$ depends on the momentum. This diffusion is described by the following expression:

$$\langle \dot{S}_{\vec{n}} \rangle_{dif} = \langle \vec{S} \frac{\delta \vec{n}}{\delta \tau} \rangle_{dif} = \langle S_{\vec{n}} \, \vec{n} \frac{\delta \vec{n}}{\delta \tau} \rangle =$$

$$= -\frac{1}{2} S_{\vec{n}} \langle \frac{(\delta \vec{n})^2}{\delta \tau} \rangle = -\frac{1}{2} S_{\vec{n}} \sum_{\alpha,\beta=1}^{3} \langle \frac{\partial \vec{n}}{\partial P_\alpha} \frac{\partial \vec{n}}{\partial P_\beta} \overline{\frac{\delta P_\alpha \delta P_\beta}{\delta \tau}} \rangle . \qquad (8)$$

The diffusion coefficient of a particle momentum may be derived by means of the expression for Tomson scattering cross-section of the photon $d\sigma_T$:

$$\overline{\frac{\delta P_\alpha \delta P_\beta}{\delta \tau}} = \delta_{\alpha\beta} \dot{N}_W \int (\delta P_\alpha)^2 d\sigma_T ,$$

where $\delta_{\alpha\beta}$ is Kroneker symbol.

As a result, we get

$$\frac{1}{2} \overline{\frac{(\delta \gamma)^2}{\delta \tau}} = \frac{56}{15} \hbar q_0^3 \omega_w \gamma^4 H_w^2 ,$$

$$\frac{1}{2} \overline{\frac{(\delta P_1)^2}{\delta \tau}} = \frac{1}{2} \overline{\frac{(\delta P_2)^2}{\delta \tau}} = \frac{4}{5} m^2 \hbar q_0^3 \omega_w \gamma^2 H_w^2 .$$

Substitution of these values in (8) yields an expression coinciding with the part α_+ quadratically dependent on the gradients of $\vec{n}$.

The last term in (1) takes into account the correlation of varying the precession axis $\vec{n}$ and the particle spin $\vec{S}$ in scattering of the photon. This correlation leads only to the spin diffusion and disappears in the classical limit. Thus, the term taking into account the correlation of changes in the spin and momentum in scattering yields an addition linear over the gradient of $\vec{n}$ to the coefficient α_+.

As seen from formulae (4), the polarization effect is altogether due to the gradients of $\vec{n}(\vec{P})$ which are determined by a field and its variables near the equilibrium orbit. This arises due to spin dependence of the radiation braking force, and its origin is purely classical, just as the coefficient $\alpha_-/\hbar$ /4/. In the classical theory the kinetic equation for $S_{\vec{n}}$, the particles with a spin $\vec{S}$, is of the form $\dot{S}_{\vec{n}} = \frac{\alpha_-}{\hbar}(\vec{S}^2 - S_{\vec{n}}^2)$ with the same coefficient α_- as in Eq.(3), and the equilibrium degree of polarization seems to be equal to 100%. Account of the quantum radiation fluctuations leads to the appearance, in the kinetic equation, of a diffusion term $-\alpha_+ S_{\vec{n}}$ decreasing, in a general case, the equilibrium degree of polarization. For the particle with a spin $\hbar/2$ the equation is converted into (3).

Let us compare the results of the present paper and Refs. /5, 6/. In /5/ the relaxation time and the degree of equilibrium polarization have been obtained on the basis of the spin-flip probabilities with respect to the incident wave direction in an accompanying system ($\zeta_\infty = 5/8$). Relative to the general formula (4), such a calculation actually assumes a zero gradi-

dient of $\vec{n}_a(\bar{P})$ in the accompanying system, for that, as seen from (6), the vector $\vec{n}(\vec{P})$ should satisfy the conditions ($\vec{n}_s = \vec{v}/v$):

$$\frac{\partial \vec{n}}{\partial P_1} = -\frac{1}{m}\vec{e}_1 , \quad \frac{\partial \vec{n}}{\partial P_2} = -\frac{1}{m}\vec{e}_2 , \quad \frac{\partial \vec{n}}{\partial P_3} = 0 .$$

The fulfilment of these conditions is possible in the case of special structure of a magnetic field of the storage ring (outside the interaction region with a wave). There is no difficulty to see that formula (4) yields the value 5/8 for ζ in this particular case.

Note, that the ratio of the spin-flip probabilities to any constant direction is not a relativistic invariant (as is the polarization direction itself) because of Thomas revolution during a momentum rotation as a result of scattering

$$\Delta\vec{\zeta}_\ell - \Delta\vec{\zeta}_a = \frac{\gamma+1}{\gamma m}\left[\vec{\zeta}\left[\vec{v}\,\Delta\vec{P}\right]\right].$$

Therefore, the probabilities given in /5/ do not characterize directly the polarization variations in the laboratory system. Indeed, a direct calculation in the laboratory system /6/ or a scale of the polarization variation from the accompanying system with account of Thomas revolution results in the equal flip probabilities (in the ultrarelativism) with respect to the incident wave direction.

An invariant quantity used in this paper is the spin projection on any space direction satisfying the equation of spin motion in an external field. In the wave-interaction region (where the storage ring field is absent) any vector constant in the electron rest system, in particular, the vector $\vec{n}(\vec{P},\tau)$ due to conditions of motion in a storage ring[3]) may be such a direction (the polarization transverse to $\vec{n}$ disappears because of spread of the spin precession frequencies in a beam).

Consider now the dependence of the equilibrium polarization degree and the relaxation time from a parameter of the orbit-spin coupling ($\rho\partial\vec{n}/\partial\rho_\perp$). Since the electrons in a wave field are mainly studied at an angle of the order of $1/\gamma$, then the degree of polarization and the characteristic relaxation time much more strongly depends on the gradient of $\vec{n}$ over the momentum in the longitudinal direction (on $\partial\vec{n}/\partial\gamma$) than on the gradients of $\vec{n}$ in the transverse direction (on $\partial\vec{n}/\partial P_{1,2}$). The dependence of the direction $\vec{n}$ from a momentum substantially increases near spin resonances. In a non-resonant situation the dependence of the values $\alpha_\pm$ from $\partial\vec{n}/\partial P_{1,2}$ may be obviously ignored $(|m\,\partial\vec{n}/\partial P_{1,2}| \ll 1)$:

$$\alpha_- = -\frac{16}{3}\hbar q_0^4 \xi_2\, \omega_w H_w^2 \gamma^3 \left(\vec{v}\frac{\partial\vec{n}}{\partial\gamma}\right)\frac{\ell}{L},$$

$$\alpha_+ = \frac{16}{3}\hbar q_0^4 \omega_w H_w^2 \gamma^2 \left\{1 - \frac{2}{5}(\vec{n}\vec{v})^2 + \frac{7}{10}\left(\gamma\frac{\partial\vec{n}}{\partial\gamma}\right)^2\right\}\frac{\ell}{L}. \qquad (9)$$

In this case, the maximum polarization degree for a circularly-polarized wave (ξ_2 = ± 1) is $\zeta_\infty = \sqrt{\frac{5}{14}} \approx 60\%$ and is achieved in a situation when the equilibrium polarization direction $\vec{n}$ in the interaction region is transverse to velocity ($\vec{n}\vec{v}$ = 0), and the gradient $\gamma\partial\vec{n}/\partial\gamma$ is directed along the velocity and equals $|\gamma\partial\vec{n}/\partial\gamma| = \sqrt{10/7}$. The relaxation time is here equal to $\left(\frac{32}{3}\hbar q_0^4 \omega_w \gamma^2 H_w^2 \frac{\ell}{L}\right)^{-1}$

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3) In the given conditions $q_0 H_w \ll \omega_w$, dynamical oscillations $\vec{n}$ under the wave field influence may be negligible.

Near the spin resonances the values $(m\partial\vec{n}/\partial P_{1,2})$ may be of the order of unity and give a significant contribution to the quantities $\alpha_\pm$. Particularly, at $\partial\vec{n}/\partial\gamma$ = 0, a maximum degree of polarization ($\approx$ 68%) is achieved in the longitudinal direction $\vec{n}$ ($\vec{n}\vec{v}$ = ± 1) in the interaction region for the gradient $\vec{n}$ equal to $m\,\partial\vec{n}/\partial P_{1,2} = \mp\sqrt{2}\,\vec{e}_{1,2}$ $(|\vec{e}_{1,2}| = 1)$. The polarization time here equals[4] $\tau_w = (5.6\hbar q_0^4 \omega_w \gamma^2 H_w^2 \ell/L)^{-1}$.

4. ON POSSIBILITIES OF LASER POLARIZATION

Let us formulate the basic requirements for a wave source. From (4.9) it follows that three characteristics of a colliding electromagnetic wave are important for the polarization process: frequency ω_w , squared field H_w^2 , which both determine the polarization time, and also the circular polarization degree of a wave ξ_2 which determines the degree of equilibrium beam polarization. No requirements for a spatial coherence are imposed $(q_0 H_w \ll \omega_w)$. There is no difficulty to generalize (4.9) in the case of polychromatic radiation since the waves with different frequencies give independent contributions to the coefficients $\alpha_\pm$.

In the polarization process a wave power is not practically supplied. This allows to increase the efficiency of a source (to increase H_w^2) by means of optical resonators and the wave focusing on the beam whose transverse sizes in a storage ring are very small (for example, in the storage ring VEPP-2M (Novo-

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4) One assumes that the spread in the spin precession frequencies remains, as before, much more less than the detuning. In the case when $m\,|\partial\vec{n}/\partial P_\perp| \approx 1$, the latter is small enough.

sibirsk) at a maximum energy is approximately equal to 10^{-4} cm^2 in the collision region). Affecting the electrons in the storage ring by the electromagnetic wave, it is necessary to take care of that the beam lifetime should be larger than the polarization time. This causes two limitations on the wave amplitude and frequency. First, in order to not knocking out the electrons in a single scattering, the wave frequency should be small enough: $4\gamma^2 \hbar \omega_W < \Delta E_p$. Second, multiple scattering leads to the particle energy diffusion in the storage ring and, hence, to increasing the energy spread. There exists a simple relation between the polarization time τ_W and the energy diffusion coefficient in the wave. For example, for an optimum value $|\gamma \partial \vec{n}/\partial \gamma| = \sqrt{10/7}$ in the storage ring we get $\tau_W^{-1} = \frac{10}{7} \frac{d}{dt} \left(\frac{\delta\gamma}{\gamma}\right)^2$. An analogous relation also occurs for the polarization process in an azimuthally-symmetric storage ring in synchrotron radiation:

$$\tau_{S.R.}^{-1} = \frac{9}{11} \frac{d}{dt} \left(\frac{\delta\gamma}{\gamma}\right)^2_{S.R.}$$

Thus, the decrease in the polarization time by a factor of Q, in comparison to a conventional storage ring ($Q = \tau_{S.R}/\tau_W$), results in increasing the energy spread by a factor of $\sqrt{1 + \frac{63}{110} Q}$. For example, the decrease in the polarization time by a factor of 10 leads to increasing the energy spread by a factor of 2.6. This increase in the energy spread does not allow to decrease considerably the polarization time[5] because of aperture limitations (size of the storage ring chamber, energy aperture).

For the best use of a source it is evidently necessary that the transverse sizes of a light beam do not exceed the established ones of the electron beam. Here, if the interaction of the beam with the wave happens many times during the radiation damping, then the polarization time decreases with a growth of the source power W inversely proportional to not the first degree of W , but only to $\sqrt{W}$, due to increasing the electron beam sizes.

A laser may be used as a source of the circularly-polarized wave. The features of laser radiation allow to focus it up to the electron beam sizes and, thereby, to increase substantially the intensity of a light field acting upon the particles. The operation in a pulse run with the phase ajusted to that of electron beam in a storage ring enables one to increase additionally the radiation intensity by a factor of L/ℓ (at the same mean power).

Moreover, one can increase the efficiency of a source, if the time between the pulses is of the order of several times of radiation beam damping in a storage ring[6]. In this case, each subsequent pulse of light fells within the already damping beam with the sizes determined only by quantum fluctuations of synchrotron radiation. Note that in the case when the energy of secondary quanta $\hbar\omega_{ph}$ is somewhat lower than the energy aperture ΔE_p , the energy loss by the particle during its interaction with a pulse of the light unessentially exceeds the energy spread arising at this time. Therefore, the appearance of the average energy deviation during interaction with the light does not decreases practically the efficiency of the considered method.

5) The increase of the energy diffusion coefficient in storage rings (at a fixed structure of the magnetic system) leads to a directly proportional increase in the beam sizes in the considered conditions when during the radiation damping the beam interacts many times with a colliding wave.

6) For example, on the storage ring VEPP-2M at a 0.5 GeV energy the time of radiation damping is $1.3 \cdot 10^{-2}$ sec.

The free-electron laser (FEL) is a promising radiation source /9/. This laser generates directly the circularly-polarized light, that enables one to use the radiation accumulated in the optical resonator of the laser for the polarization process. The ondulator and the mirrors are here installed directly in a straight section of the storage ring. Since the electrons with an energy of several tens of MeV play a role of the active substance, their passing through the optical resonator mounted in the storage ring do not worsen the conditions necessary for keeping the circular beam.

For illustration, let us present a numerical example. In conventional storage rings the spin precesses around the vertical direction and a quantity $|\gamma \partial \vec{n}/\partial \gamma|$ is connected only to the perturbations of magnetic system[7]. One of the possible methods of organizing the gradient $\gamma \partial \vec{n}/\partial \gamma$ directed along $\vec{v}$ on the wave-interaction region is the creation of a longitudinal field in the opposite straight section of the storage ring (solenoid). For the optimum case, when $\gamma \partial \vec{n}/\partial \gamma = \sqrt{\frac{10}{7}}\,\vec{v}$, the solenoid is required whose length ℓ_{sol} and field H_{sol} satisfy the following relation[8]: $H_{sol}\ell_{sol} \approx 10(\nu-K)^2\,[T\times M]$, where ν is the spin-precession frequency in the storage ring in the units of revolution frequency, and K is the number of the nearest integer spin resonance. It should be noted that at a constant detuning from the resonance the needed integral of the solenoid field is independent of the particle energy in the storage ring. The growth of depolarising processes assosiated with imperfection of a magnetic system of the storage ring does not allow to choose the values ν as close to the integer as possible[9]. In the case of using the laser with a mean power W in the optimum case (i.e. in the case of a pulse régime with the phase ajusted to that of electron beam, and a correspondence between the light beam sizes and the particle beam sizes) the polarization time may be presented in the form:

$$\tau_w^{-1}(сек^{-1}) = 1.2\cdot 10^{-13}\,\frac{E^2[GeV^2]\;W[w]}{S[см^2]\;\lambda_w[см]},$$

where S is the particle beam area.

For the laser on Yt-Al granate with Nd with the following parameters: $W = 10^3$ W, $\lambda_w = 10^{-4}$ cm /10/, the polarization time equals 330 sec at 0.5 GeV and with a beam area 10^{-4} cm^2.

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7) During the passing of the section with a wave (helical) field, strictly speaking, the vector $\vec{n}$ is inverted around the velocity by an angle $[(q'/q_0)^2+\gamma^{-2}](q_0 H_w/\omega_w)^2\omega_w \ell/2$. This angle is usually small and may be ignored.

8) The written relation is valid at $|\nu-K| \ll 1$.

9) For example, for the electron polarization by a colliding wave at 10 minutes in the storage ring VEPP-2M (Novosibirsk) it is necessary to choose the detuning values $|\nu-K|$ larger than 10^{-2}.

R e f e r e n c e s

1. A.N.Skrinsky. Report at XVIII International Conference on High Energy Physics. Tbilisi, July 1976.
2. Ya.S.Derbenev, A.M.Kondratenko, S.I.Serednyakov, A.N.Skrinsky, G.M.Tumaikin, Yu.M.Shatunov. Particle Accelerators **8** (1978).
3. A.Hutton. Part. Acc. **7**, 177 (1976).
4. Ya.S.Derbenev, A.M.Kondratenko. ZhETF, 64, 1918 (1973).
5. Yu.A.Bashmakov, E.G.Bessonov. Proceedings of All-Union Meeting on Charged Particle Accelerators. "Nauka", Moscow 1977, vol.I, p.277.
6. Yu.G.Pavlenko, A.Kh.Mussa 'Vestnik Moskovskogo universiteta', fiz. astr. 18, N°2, 59 (1977).
7. Ya.S.Derbenev, A.M.Kondratenko, A.N.Skrinsky. 'Reports of the Academy of Sciences of the USSR', 192, 1255 (1970).
8. F.Lipps, H.A.Tolhock. 'Physica', 20, 395 (1954).
9. D.A.G.Deacon, L.R.Elias, J.M.J.Madey, G.J.Ramian, H.A.Schwettman, T.I.Smith. 'Phys.Rev.Lett,', 38, N°16, 892 (1977).
10. J.Ready. 'Effects of high-power laser radiation', Academic Press, New-York-London, 1971.

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